The $b\to s\gamma$ constraint in effective supergravities from string theory

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ABSTRACT

We study the constraints from the $b\to s\gamma$ decay in the parameter space of effective supergravities from orbifold string theory and with minimal supesymmetric particle content. Both the general dilaton-dominated universal scenario as well as a non-universal scenario for the soft terms are investigated. It is found that the recently reported CLEO upper and lower bounds constrain the parameter space of the models under scrutiny. In particular we find constraints on the values of the parameter $\tan\beta$ and the gluino masses. In this class of string scenarios the negative sign of the Higgs mixing parameter μ, is phenomenologically preferred.

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1   Introduction

One of the prime tasks of the Large Hadron Collider is to search for the supersymmetric partners of the Standard Model multiplets. Once the first sparticles are discovered, the program of sparticle spectroscopy will give us vital clues for the *underlying* theory that explains the observed spectrum. As is well known $N = 1$ superymmetric field theories predict a very rich structure of sparticles from a few fundamental parameters. These parameters break the supersymmetry softly at an energy of the order of the electroweak scale thus ensuring that the hierarchy problem is at least technically solved and that the superparticles do not have the same mass as their Standard Model partners. It is customary to parametrize the effects of the soft-supersymmetry breaking terms by four universal parameters: the universal gaugino mass $M_{1/2}$, the scalar terms associated with the trilinear couplings in the superpotential A, the scalar masses m_0, and the B term associated with the Higgs-doublet mixing term in the superpotential [2].

On the other hand, in order to be able to interpret the "lines" in the sparticle spectrum we need to have a theory which will differentiate among the many alternatives of the soft-parameter space. The only example of the sort of theory we are aiming for is heterotic string theory. In string theory these soft susy- breaking parameters are in principle, calculable, but a definite answer is at present lacking due to the fact that the superymmetry breaking mechanism in the theory is not well understood. However, in the pioneering work of [1] the effect of SUSY-breaking is parametrized by the VEVs of the F-terms of the dilaton (S) and the moduli (T_m) chiral superfields, generically present in large classes of four-dimensional supersymmetric heterotic strings. This is an important step towards a theory which will explain the rich sparticle spectrum. In this work the soft-parameter space has been reduced since many interesting relations among the soft-parameters have been found which in principle can be tested at the LHC .

However, until the first experiments at LHC start running, we have to use all the current experimental information in order to study the parameter space of the effective supergravities from string theory. Unfortunately, most of the precision LEP measurements are not very sensitive to new physics as the Standard Model contributions enter at the tree level, while possible new physics contributions begin at the one loop level. Thus the most one might hope for in these measurements is a few percent correction from new physics.

However, it has become well known that the $b{\rightarrow}s\gamma$ decay is an exception to this and that is a powerful tool for testing Physics beyond the Standard Model [2, 3]. This is the case for several reasons: First as a FCNC process, the $b{\rightarrow}s\gamma$ decay, arises first at the one loop level so that the Standard Model loops and new physics loops enter at the same level. Second, the decay is of size $G_F^2\alpha$, where G_F is the Fermi constant (rather then $G_F^2\alpha^2$ as is usual for FCNC processes). Third,

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[2]Note however, that this parametrization is not the most general case since in string theory non-universality of the soft terms is very common [1].

there are available experimental data on the exclusive [4] $B{\to}K^*\gamma$ and the inclusive [5] $B{\to}X_s\gamma$ decays that lead to upper and lower bounds on the branching ratio BR($b{\to}s\gamma$) of the same order as the SM prediction. In particular, from the inclusive B decay:$1{\times}10^{-4}{<}$BR($b{\to}s\gamma$)$<4{\times}10^{-4}$ [6]. In the Minimal Supersymmetric extension of the Standard Model there are additional contributions to the decay besides the SM diagram with a W gauge boson and a top quark in the loop. In particular, there are additional contributions coming from loops involving charged Higgses (H^-) and a top quark, charginos (χ^-) and u-type squarks (of which the relevant contributions come from the stops, $t_{L,R}$, and scharms, $\tilde{c}$, and a gluino or neutralinos (χ_i^0) plus a d-type squark (mainly $\tilde{b}$ and $\tilde{s}$) [3]. As pointed in ref. [3], the latter two diagrams do not contribute significantly to the BR and can therefore be neglected. It is common practice to use the ratio defined as

$$R = \frac{BR(b{\to}s\gamma)}{BR(b{\to}ce\bar{\nu})} \simeq \frac{BR(B{\to}X_s\gamma)}{BR(B{\to}X_ce\bar{\nu}_e)} \tag{1}$$

to constrain various models, utilizing the well determined value of $10.7 \pm 0.5\%$ for $BR(B{\to}X_ce\bar{\nu}_e)$. The advantage of using R, instead of $BR(b{\to}s\gamma)$, is that the latter is dependent upon m_b^5 while the former only depends on $z = m_c/m_b$, the ratio between the c and b quark masses, which is much better determined than both masses, i.e., $z = 0.316 \pm 0.013$ [7].

The ratio R defined in Eq. (1) is given by [3]

$$R = \frac{|V_{ts}^*V_{tb}|^2}{|V_{cb}|^2} \frac{6\alpha_{QED}}{\pi} \frac{[\eta^{16/23}A_\gamma + \frac{8}{3}(\eta^{14/23} - \eta^{16/23})A_g + C]^2}{I(m_c/m_b)[1 - (2/3\pi)\alpha_s(m_b)f(m_c/m_b)]} \tag{2}$$

where $\eta = \alpha_s(M_W)/\alpha_s(m_b)$, and M_W is the W boson mass. Here, $I(z) = 1-8z^2+8z^6-z^8 - 24z^4 \ln z$ is the phase-space factor, and $f(z) = 2.41$, is a QCD correction factor, for the semileptonic process, $b{\to}ce\bar{\nu}_e$. C represents the leading-order QCD corrections to the $b{\to}s\gamma$ amplitude when evaluated at the $Q = m_b$ scale [8] [3]. In evaluating Eq. (2) we also take $|V_{ts}^*V_{tb}|^2/|V_{cb}|^2 = 0.95 \pm 0.04$. Finally, $A_{\gamma,g}$ are the coefficients of the effective operators for the $bs\gamma$ and bsg interactions; in our case, as mentioned above we consider as relevant the contributions coming from the SM diagram plus those with top quark and charged Higgs, and stops/scharms and charginos running in the loop. Their expressions are given by:

$$A_{\gamma,g}^{SM} = \frac{3}{2}\frac{m_t^2}{M_W^2}f_{\gamma,g}^{(1)}\left(\frac{m_t^2}{M_W^2}\right)$$

$$A_{\gamma,g}^{H^-} = \frac{1}{2}\frac{m_t^2}{m_H^2}\left[\frac{1}{\tan^2\beta}f_{\gamma,g}^{(1)}\left(\frac{m_t^2}{m_H^2}\right) + f_{\gamma,g}^{(2)}\left(\frac{m_t^2}{m_H^2}\right)\right]$$

$$A_{\gamma,g}^{\chi^-} = \sum_{j=1}^2\left\{\frac{M_W^2}{M_{\chi_j}^2}\left[|V_{j1}|^2 f_{\gamma,g}^{(1)}\left(\frac{m_{\tilde{c}}^2}{M_{\chi_j}^2}\right) - \sum_{k=1}^2\left|V_{j1}T_{k1} - \frac{V_{j2}m_tT_{k2}}{\sqrt{2}M_W\sin\beta}\right|^2\right.\right.$$

[3]For a discussion about the uncertainties derived from the selection of renormalization at the $Q = m_b$ scale see [9].

2

$$\times \quad f^{(1)}_{\gamma,g}\left(\frac{m^2_{\tilde{t}_k}}{M^2_{\chi_j}}\right)\Bigg] - \frac{U_{j2}}{\sqrt{2}\cos\beta}\frac{M_W}{M_{\chi_j}}\left[V_{j1}f^3_{\gamma,g}\left(\frac{m^2_{\tilde{c}}}{M^2_{\chi_j}}\right)\right.$$

$$- \quad \sum_{k=1}^{2}\left(V_{j1}T_{k1} - V_{j2}T_{k2}\frac{m_t}{\sqrt{2}M_W\sin\beta}\right)T_{k1}f^{(3)}_{\gamma,g}\left(\frac{m^2_{\tilde{t}_k}}{M^2_{\chi_j}}\right)\Bigg]\Bigg\} \tag{3}$$

where the functions $f^i_{\gamma,g}$, $i = 1,2,3$ may be found in [3], and all the masses are understood to be at the electroweak scale. V and U are the matrices which diagonalise the chargino mass matrix, while T diagonalises the stop mass matrix.

Our strategy is to use the renormalization group equations (RGEs) to calculate the mass spectrum subject to combined constraints from experiment and correct radiative electroweak breaking, using as boundary conditions for the soft-susy breaking terms string scenarios obtained in ref. [1]. We then use this spectrum to evaluate the ratio R via Eq. (1) and Eq. (2) to obtain

$$BR(B \to X_s\gamma) = R \times BR(B \to X_c e\bar{\nu}_e) \tag{4}$$

and compare the results with the CLEO II bound. Let us describe now our renormalization group procedure in detail. We want of course to integrate our RGEs from the string unification scale M_{string} of $O(10^{16})$GeV down to the electroweak scale M_Z [4] . In order that our numerical integration routines make the run starting from M_{string}, we must have the values of all the parameters at this scale, but this is difficult to achieve. The problem is that the values of many parameters are known experimentally at low scales. However, the values of other parameters, such as soft breaking terms, are most easily understood at higher energies where theoretical simplification may be invoked. Thus, there is no scale at which there is both theoretical simplicity and experimental data. In other words, choosing M_{string} as our starting point we want to find the M_{string} values of all the parameters such that we recover the expected low energy values after renormalization group evolution to experimental scales. An approach that incorporates some boundary conditions at both electroweak and M_{string} scales, is the so called ambidextrous approach [11, 10]. In this approach one specifies m_t and $\tan\beta$ at the electroweak scale (along with M_Z and M_W) and $M_{1/2}$, m_0, and A at the M_{string} scale. Thus first one integrates the dimensionless parameters given $\tan\beta(M_Z)$ and $m_t(M_Z)$ up to M_{string} in order to specify the complete set of boundary conditions at this scale. Then all the parameters including the soft SUSY-breaking are evolved from M_{string} to the electroweak scale. At this scale $\mu(M_Z)$ and $B(M_Z)$ are determined by minimizing the one-loop effective potential V^{1-loop} . Subsequently μ and B can be RGE-evolved up to the string scale. This strategy is effective because the RGEs for the soft-supersymmetry breaking parameters do not depend on μ and B. This method has two powerful advantages. First, any point in the $m_t - \tan\beta$ plane can be readily investigated in specific supergravity models since m_t and $\tan\beta$ are taken as inputs. This is extremely useful since after the recent discovery of the top quark at

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[4]we assume that string threshold corrections are such that they explain the mismatch between the unification scale and the string scale [1].

Fermilab its mass is going to be determined with high accuracy in the future. Thus approaches like the top-down approach in which m_t is output cannot effectively scan the parameter space of supergravity models. Secondly, the minimization conditions of V^{1-loop} are easily solved for μ and B. Specifically, we employ a two-dimensional Newton method from the NAG library which quickly locates the extremal values for μ, B by iteration.

2 Effective Supergravities from String Theory.

The low-energy limit of the supertring models relevant for the phenomenology is the $N = 1$ supergravity (SUGRA) described by the Kähler function G, which is a function of the Kähler potential K and the superpotential W, and the gauge kinetic functions f_a [12]. The generic fields present in the massless string spectrum contain the dilaton superfield S, moduli fields generically denoted by T_i (which can contain the radii-type moduli T_i and the complex structure moduli U_j) and some matter chiral fields ϕ^α, containing the Standard Model particles. The resulting effective low-energy theory, emerging from string theory, possesses a high degree of symmetry, which in general restricts the form of the three SUGRA functions mentioned above. As a result, the soft parameters are also constrained [5].

A particularly interesting class of such *stringy* symmetries are the *target-space duality* symmetries. The physical content of such symmetries is that in string theory physics at a very small scale cannot be distinguished from physics at a very large scale. Under such symmetries the moduli fields T_i transform as

$$T_i \rightarrow \frac{a_i T_i - i b_i}{i c_i T_i + d_i} \ , a_i d_i - b_i c_i = 1 \ , a_i \ldots d_i \in Z. \tag{5}$$

In effective string theories of the orbifold type [13], the matter fields ϕ^α transform under (5) as

$$\phi^\alpha \rightarrow (i c_i T_i + d_i)^{n_\alpha^{(i)}} \phi^\alpha \tag{6}$$

where the integers $n_\alpha^{(i)}$ are called modular weights (usually are negative integers). With the above transformations the Kähler G function is modular invariant (if the superpotential W has modular weight -3).

The Kähler potential K (to first order in the observable fields) is given in general by the form [1, 14]

$$K = -\log(S + S^*) + K_0(T, T^*) + K_{\alpha\beta}(T, T^*)\phi_\alpha \phi_\beta^* \tag{7}$$

where the indices α, β label the charged matter fields. The authors in ref. (1) concentrated in the case of the overall modulus T and disregarded any mixing between the S and T fields kinetic terms which is strictly correct at the tree level. At one loop level such a mixing arises through the Green-Schwarz mixing coefficient δ_{GS}^i [15].

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[5]This is to be contrasted with conventional SUGRA theories where G and f are arbitrary

The scalar potential in the low-energy supergravity action has the form [12]

$$V = |W_{SUSY-breaking}(T,S)|^2 e^{K_0}(G^i (G^{-1})^j_i G_j - 3) \qquad (8)$$

($e^G = |W|^2 e^K, G_i = \partial G/\partial \phi_i$.) In deriving (8) the authors in [1] assumed that, upon minimization of V,$\langle G_\alpha \rangle = 0$ and $\langle Q_\alpha \rangle = 0$ in the matter sector. This assumption, which is satisfied in most realistic scenarios, means that the spontaneous supersymmety breaking takes place in the dilaton-moduli sector, i.e. $\langle G_i \rangle \neq 0$ for at least one of the moduli fields. Then the gravitino mass becomes

$$m_{3/2} = e^{K_0(T,S,T^*,S^*)/2}|W_{SUSY-breaking}(T,S)| \qquad (9)$$

$m_{3/2}$ should be of order TeV. Then one can obtains the following soft terms;first the gaugino masses take the form

$$M_a(T,T^*,S,S^*) = \frac{1}{2}m_{3/2}G^i(T,S,T^*,S^*)\partial_i \log g_a^2(T,S,T^*,S^*) \qquad (10)$$

The scalar masses (squarks and sleptons) become [1]

$$m^2_{\alpha\bar{\beta}}[K_{\alpha\bar{\beta}}(T,S,S^*,T^*) - G^i(T,S,S^*,T^*)G^{\bar{j}}(T,S,S^*,T^*)R_{i\bar{j}\alpha\bar{\beta}}] \qquad (11)$$

($R_{i\bar{j}\alpha\bar{\beta}} = \partial_i\bar{\partial}_{\bar{j}}K_{\alpha\bar{\beta}} - \Gamma^\gamma_{i\alpha}K_{\gamma\delta}\bar{\Gamma}^{\bar{\delta}}_{\bar{j}\bar{\beta}}, \Gamma^\gamma_{i\alpha} = K^{\gamma\bar{\delta}}\partial_i K_{\alpha\bar{\gamma}}$)

By assuming that SUSY breaking is triggered by the auxiliary fields of the dilaton-moduli sector, one can parametrize the unknown supersymmetry dynamics by some angle $\tan\theta = \langle F_S \rangle / \langle F_T \rangle$ [1]. Then the exact form of the (perturbative or non-perturbative) superpotential is parametrized by θ and $m_{3/2}$, and the form of the soft-parameters depend only on known perturbative quantities like K. Next we discuss the different scenarios that emerge in this framework which are subject of our research in this paper.

3 Models

Using the general expressions (10), (11) the following form of soft terms may be derived [1] [6]

$$m^2_\alpha = m^2_{3/2}[1 + n_\alpha \cos^2\theta] \qquad (12)$$

$$M_a = \sqrt{3}m_{3/2}\frac{k_a ReS}{Ref_a}\sin\theta + m_{3/2}\cos\theta\frac{B'_a(T+T^*)\hat{G}_2(T,T^*)}{32\pi^3 Ref_a} \qquad (13)$$

$$A_{\alpha\beta\gamma} = -\sqrt{3}m_{3/2}\sin\theta - m_{3/2}\cos\theta(3 + n_\alpha + n_\beta + n_\gamma) \qquad (14)$$

[6]we consider no-scale scenarios in which the cosmological constant is zero. If one relaxes the constraint of the vanishing of the cosmological constant the soft terms depend explicitly on the nonzero value of the latter and consequently the experimental predictions of the models [1].

where k_a is the Kac-Moody level of the gauge factor. In the phenomenological analysis that follows $k_3 = k_2 = \frac{3}{5}k_1 = 1$ and the definitions of B'_a, $\hat{G}_2$ functions may be found in [1].

As regards the B soft term associated with the Higgs mixing μ term in the superpotential its form is model dependent. In particular its value depends on the scenario we use for the generation of the μ term [7]. In string theory we have

• The quadratic μ term arises as an effective non- renormalizable fourth- (or higher) order term in the superpotential of the form

$$\lambda W_0 H_1 H_2 \tag{15}$$

where W_0 is the renormalizable superpotential and λ an unknown coupling, which mixes the observable sector with the hidden sector, then a μ term is automatically generated with size $\mu = \lambda m_{3/2}$ [17].

• The quadratic μ term is built into the theory through the Kähler potential, and becomes non-zero and of $O(m_{3/2})$ upon superymmetry breaking [18, ?, 20].

In no-scale scenarios the value of B in both cases is given by

$$B = 2m_{3/2} \tag{16}$$

As one can see the soft terms are in general $non - universal$. However for $\theta = \frac{\pi}{2}$, i.e the dilaton dominated supersymmetry breaking and neglecting threshold corrections, the soft terms are in fact universal [1]:

$$m_0 = \frac{1}{\sqrt{3}}M_{1/2}, \quad A = -M_{1/2} \tag{17}$$

In the strict dilaton-dominated scenario the B soft term is also predicted to have the value

$$B = 2m_0 = \frac{2}{\sqrt{3}}M_{1/2} \tag{18}$$

However, in the general dilaton-dominated scenario the B term is an independent parameter. The latter scenario is subject of our research in this paper. In the numerical approach we use the B term is determined by the minimization conditions.

The special of properties of the dilaton dominated scenario have been recently *emphasized* in ref. [21]. It is also worthwhile to reiterate that the dilaton dominated scenario is of general validity since the boundary conditions in (17) are obtained for *any* 4-D N=1 string and not only for orbifolds. An initial study of the phenomenology of the above soft terms in the context of the Minimal Supersymmetric Standard Model (MSSM) has been done in [22, 1]. In ref.[23], Lopez et al, studied the phenomenological consequences of (17) in the context of $SU(5) \times U(1)$ which predicts extra matter particles [8].

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[7]For a recent review see Ref. [16]

[8]In string models derived in the fermionic formulation the extra matter hypothesis is imperative for reconciling the string unification scale with the LEP data, at least at the $k = 1$ level [24].

4 Analysis

As was said in the introduction we use the ambidextrous approach in our RG analysis. Thus, our parameter space in the dilaton-dominance limit is $\tan\beta$, $m_t(M_Z)$, $M_{1/2}$ and the sign of μ which is not determined by the radiative electroweak breaking constraint. In the more general case where the moduli also contribute to SUSY breaking the goldstino angle is added to the parameter space. In the latter case, we must take into account additional $D-$ term contributions to the scalar masses due to the non-universality of the scalar soft-terms in (12) [25, 26]. In particular, the combination

$$
\begin{aligned}
\mathcal{S} \; = \; & m_{H_2}^2 - m_{H_1}^2 \\
+ \; & \mathrm{Tr}[\mathbf{M}_{Q_L}^2 - \mathbf{M}_{L_L}^2 - 2\mathbf{M}_{U_R}^2 + \mathbf{M}_{D_R}^2 + \mathbf{M}_{E_R}^2]
\end{aligned}
\tag{19}
$$

contributes to the RGEs and satisfies the (one-loop) scaling equation

$$
\frac{d\mathcal{S}}{dt} = \frac{2b_1 g_1^2}{16\pi^2}\mathcal{S}
\tag{20}
$$

so that if it is zero at some scale, for example the string scale, then it is zero for all scales. The renormalization group coefficient $b_1 = 33/5$ in the MSSM. In the dilaton-dominated scenario the $\mathcal{S}$ term does not contribute to the scaling of scalar masses. In the non-universal case we consider the model with modular weights n_α different from -1 which was first studied in [1] and gives unification at a scale of $O(10^{16})GeV$. Again we prefer to allow the B soft term to be a free parameter given the uncertainty conserning the μ term. However, the results as regards the $b{\rightarrow}s\gamma$ decay in the latter model are similar to those obtained in the dilaton-dominated scenario. The reason is that the $\sin\theta$ parameter takes values close to one in order to avoid tachyonic states and therefore the dilaton $F-$ term is the dominant source of supersymmetry breaking see Brignole et al in [1].

The string low energy observable sector is identified with that of the MSSM, and the perturbative superpotential which describes the renormalizable trilinear and bilinear Yukawa couplings of quarks,leptons and Higgs bosons chiral superfields is given by

$$
\begin{aligned}
W \; = \; & \sum_i h_{u_i} Q_i H_2 U_i^c + h_{d_i} Q_i H_1 D_i^c + h_{e_i} L_i H_i E_i^c \\
+ \; & \mu H_1 H_2,
\end{aligned}
\tag{21}
$$

where i is a generation index, $Q_i(L_i$ are the scalar partners of the quark (lepton) SU(2) doublets, $U_i^c, D_i^c(E_i^c)$ are the quark (lepton) singlets and $H_{1,2}$ are the two supersymmetric Higgs doublets. The $h-$ factors are the Yukawa couplings and μ is the usual Higgs mixing parameter. In Eq. (21) the usual SU(2) contraction is assumed, e.g $\mu\epsilon_i j H_1^i H_2^j$ with $\epsilon_{12} = -\epsilon_{21} = 1$. Then the Lagrangian will contain

7

besides the superymmetric $F-$ and $D-$ terms the following soft supersymmetric breaking terms

$$
\begin{aligned}
\mathcal{L}_{SB} &= \frac{1}{2}M_a\lambda_a\lambda_a + \left\{ \sum_{i,j}(m^2)^j_i\phi^i\phi_j \right. \\
&+ \sum_i [A_{u_i}h_{u_i}\tilde{Q}_i H_2\tilde{U}^c_i + A_{d_i}h_{d_i}\tilde{Q}_i H_1\tilde{D}^c_i + A_{e_i}h_{e_i}\tilde{L}_i H_1\tilde{E}^c_i + \text{h.c}] \\
&+ \left. [B\mu H_1 H_2 + \text{h.c}] \right\} \quad ,
\end{aligned}
\tag{22}
$$

where ϕ_i denotes a generic scalar field.

For the study of radiative electroweak breaking constraint (REWB) we use the one-loop effective potential instead of the tree-level potential V_0

$$
V^{1-loop} = V_0 + \Delta V_1 \tag{23}
$$

where

$$
\Delta V_1 = \frac{1}{64\pi^2}\text{Str}\left[\mathcal{M}^4\left(\log\frac{\mathcal{M}^2}{Q^2} - \frac{3}{2} \right) \right] \tag{24}
$$

depend on the Higgs fields through the tree-level squared-mass matrix $\mathcal{M}^2$. The supertrace in (24) is given by

$$
\text{Str} f(\mathcal{M}^2) = \sum_i (-1)^{2J_i}(2J_i + 1)f(m^2_i) \tag{25}
$$

where m^2_i denotes the field-dependent mass eigenvalue of the ith particle of spin J_i. For the calculation of radiative corrections we use the tadpole method [27] which is a very convienient way of incorporating the corrections into the minimization conditions of all the particle spectrum.

The chargino mass term in matrix form, which plays a crusial role in the expressions for $BR(B{\to}s\gamma)$ [see Eq.(3)] is given by

$$
\begin{pmatrix} M_2 & \sqrt{2}m_W\sin\beta \\ \sqrt{2}m_W\cos\beta & -\mu \end{pmatrix}
$$

Besides the constraint of correct electroweak breaking, the experimental constraints we impose in the above superstring scenarios are (1) We require that all sleptons be heavier than $M_Z/2$, since sleptons are not observed in Z decays [28]. (2) We require that the lightest chargino mass eigenstate, $M_{\tilde{\chi}^+_1}$, be heavier than $M_Z/2$, since chargino pairs are not observed in Z decays [28]. (3) We impose that gluinos be heavier than 120GeV. However, this requirement is not so constraining since the sleptons and chargino boundary conditions require that $M_g{>}200$ GeV. Because of naturalness criteria the largest gluino masses we study correspond to $M_g{\approx}1$TeV. (4) As regards the Higgs sector we require that the lightest Higgs eigenstate, h^0, is heavier than 60GeV and that the CP-Odd mass eigenstate A^0 is not visible at LEP. (5) We

8

demand that all squarks should be heavier than 45GeV and the lighest neutralino be heavier than 20 GeV, and $m_{top} = 178GeV$ (6) Finally, as was said in the introduction we impose the current CLEO bounds on the $BR(b{\to}s\gamma)$

As one can see from the graphs we plotted the values of $BR(b{\to}s\gamma)$ vs the gluino mass ,in the dilaton-dominated scenario and in the non-universal case with $\theta = \frac{2\pi}{3}$,for selected values of $\tan\beta$ and for the top mass $m_t^{\text{pole}} = 178GeV$ consistent with the experimental values that were announced from CDF recently [29]. The SM prediction for the BR and the CLEO bounds are also shown . From the graphs is evident that the CLEO upper and lower bounds restrict the allowed parameter space dramatically and in fact require μ to be negative. For $\mu{>}0$ the values of $BR(b{\to}s\gamma)$ increase steadily with $\tan\beta$ and fall outside the experimentally allowed region for all values of M_g and for $\tan\beta \geq 2$. Thus, we see that the upper CLEO bound together with the recently announced value for the top quark mass [29] exclude the positive branch of the Higgs mixing parameter μ. For $\mu{<}0$ the $\tan\beta-$dependence is different. One sees that $BR(b{\to}s\gamma)$ can be suppressed much below the lower CLEO bound and consequently of the Standard Model result. This phenomenon has been explained in [30, 31]. The chargino contribution to the amplitude in Eq. (3), can have the same sign (negative) or opposite sign (positive) compared to $t - W^{\pm}$ and $t - H^+$ contributions which are always negative. Actually, the region in which the chargino amplitude gives rise to a destructive interference effect with the other amplitudes corresponds to the region in which μ is negative [9] Thus, constructive interference occurs for $\mu{>}0$ and destructive interference occurs for $\mu{<}0$, as evident from the figures.

Furthermore, for the phenomenologically prefered negative branch of the μ Higgs mixing term, we observe a tendency towards smaller values of the ratio of the two vacuum expectation values. As $\tan\beta$ increases a larger portion of the parameter space is excluded and higher gluino masses are preferred. Actually $tan\beta_{MAX}{\approx}30$ in this case since higher values of this parameter which give correct electroweak breaking are very expencive and demand very high gluino masses, above the naturaleness bound of 1TeV. Here, the lower CLEO bound is relevant for the constraints described. Thus, we can conclude that the experimental evidence for the inclusive $b{\to}s\gamma$ decay together with the recent top quark discovery remain among the most relevant tests for exploring the parameter space of superstring scenarios.

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[9]Notice that the renormalization for μ is in fact multiplicative.

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Figure 1: The BR in the dilaton-dominated scenario for $\tan\beta = 2$ and $\mu < 0$

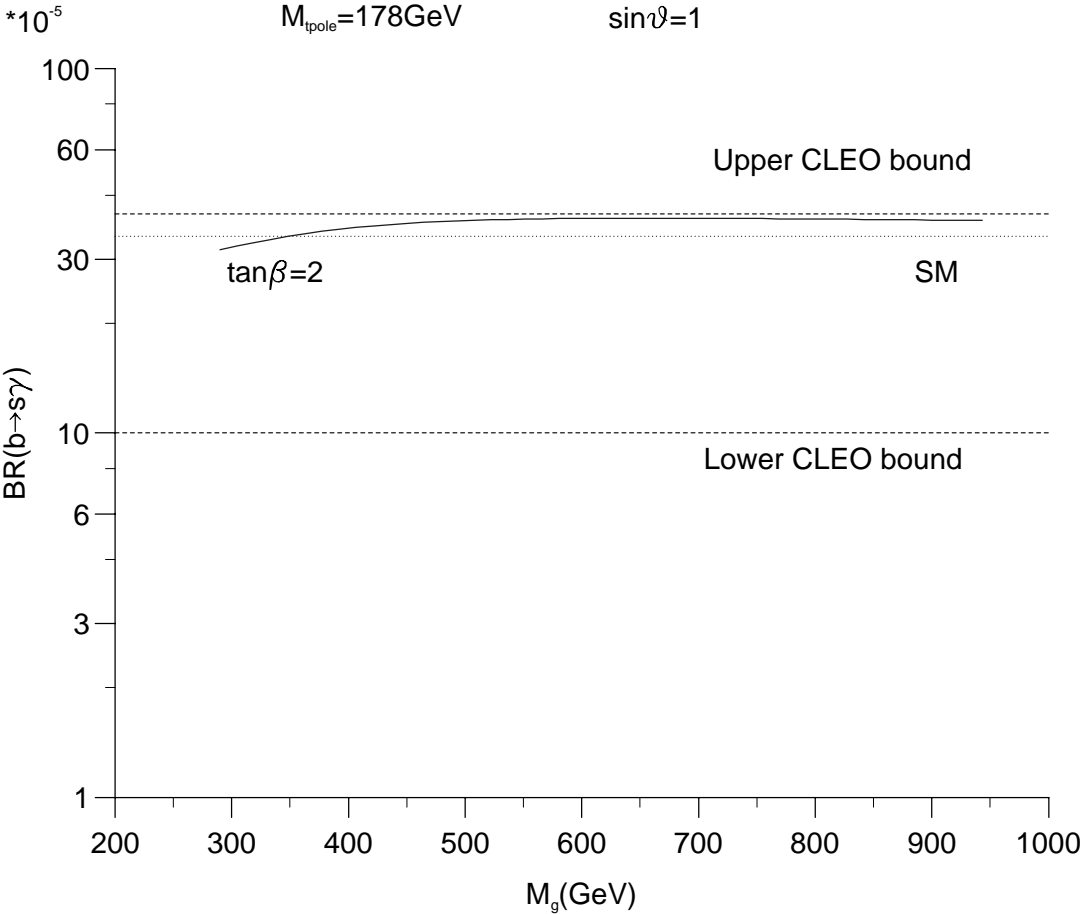Figure 2: The BR in the dilaton-dominated scenario for $\tan\beta = 6$ and $\mu < 0$

Figure 3: *The ratio in the dilaton-dominated scenario for $\mu < 0$, $\tan\beta = 10$*

Figure 4: The BR in the dilaton-dominated scenario for $\tan\beta = 15$, $M_{tpole} = 178\text{GeV}$,and $\mu < 0$

Figure 5: The BR in the dilaton-dominated scenario for $\tan\beta = 20$ and $\mu < 0$

Figure 6: The BR in the dilaton-dominated scenario for $\tan \beta = 23$ and $\mu < 0$

Figure 7: *The ratio in the dilaton-dominated scenario for $\mu > 0$ and different choices for* $\tan\beta$ $|V_{ts}^* V_{tb}|^2 / |V_{cb}|^2 = 0.99$.

Figure 8: *The ratio in the dilaton-dominated scenario for* $\tan\beta = 30, \mu{>}0$

Figure 9: *The ratio in the dilaton-dominated scenario for* $\tan\beta = 2, 9, 15$ *and* $|V_{ts}^* V_{tb}|^2 / |V_{cb}|^2 = 0.95$

Figure 10: The BR in the mixed scenario, $\theta = \frac{2\pi}{3}$, $\mu < 0$ and different choices of $\tan \beta$
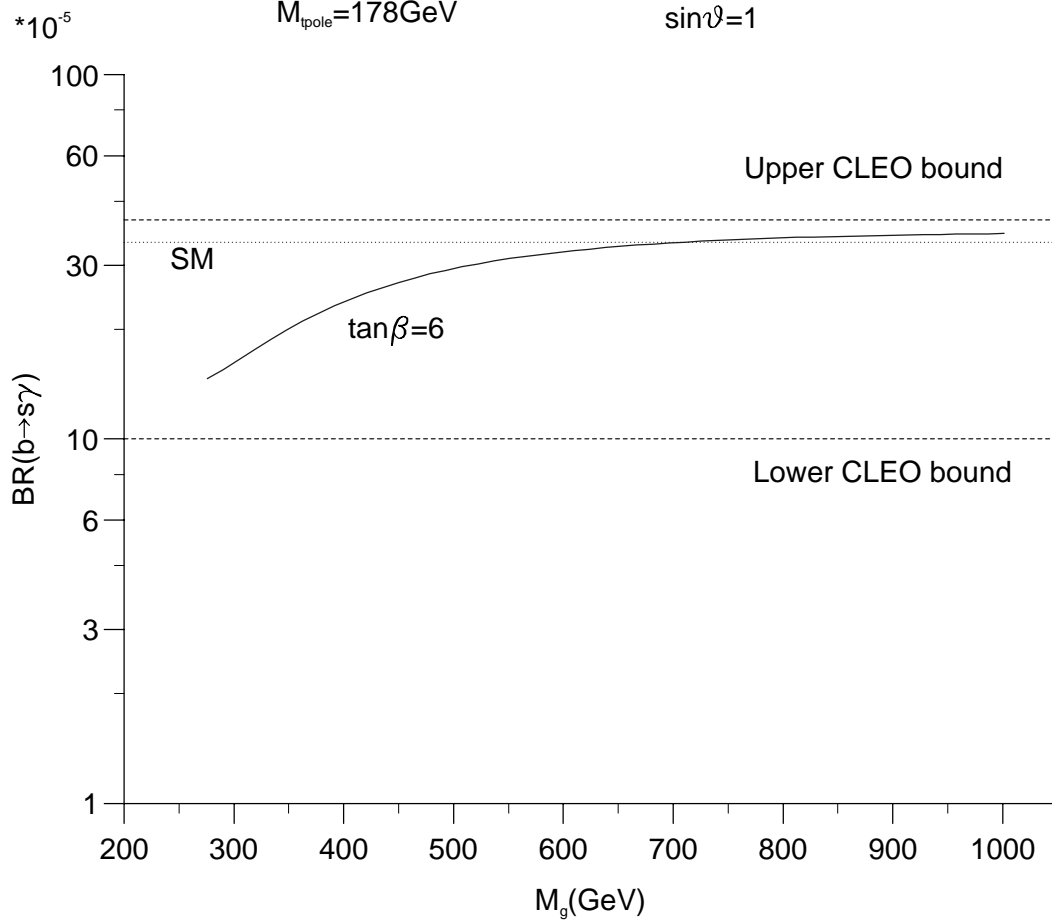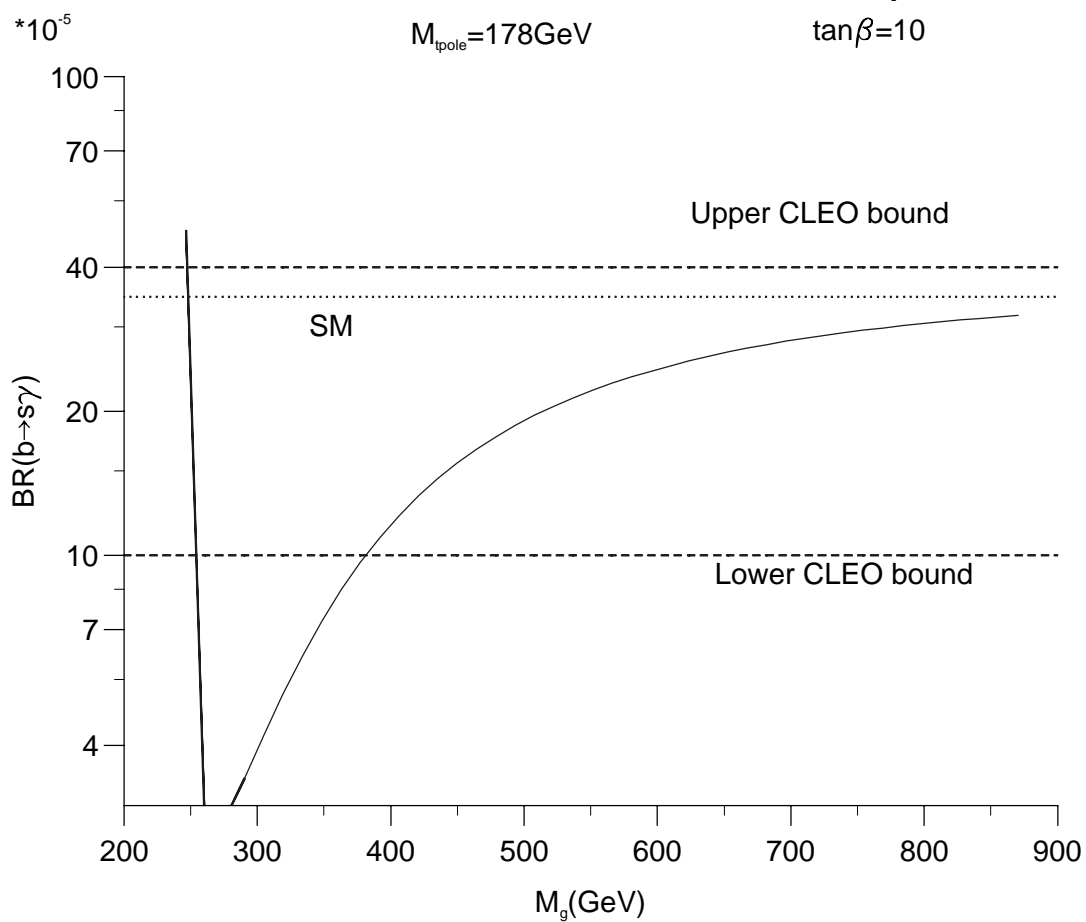
The BR vs gluino mass for $\mu<0$

$M_{tpole}=178$GeV $\sin\vartheta=1$

$*10^{-5}$

Upper CLEO bound

$\tan\beta=2$ SM

Lower CLEO bound

BR($b{\to}s\gamma$)

M_g(GeV)

The BR vs gluino mass for $\mu<0$

$M_{tpole}=178\,GeV$ $\sin\vartheta=1$

$*10^{-5}$

Upper CLEO bound

SM

$\tan\beta=6$

Lower CLEO bound

BR(b→sγ)

$M_g(GeV)$

The BR vs gluino mass for $\mu<0$

$*10^{-5}$

$M_{tpole}=178\text{GeV}$　　　　$\tan\beta=10$

Upper CLEO bound

SM

Lower CLEO bound

BR(b→sγ)

$M_g(\text{GeV})$

The BR vs gluino mass for $\mu<0$

$*10^{-5}$

$M_{tpole}=178GeV$

Upper CLEO bound

SM

$tan\beta=15$

BR($b\rightarrow s\gamma$)

Lower CLEO bound

M_g(GeV)

The BR vs gluino mass for $\mu<0$

$M_{tpole}=178$GeV $\sin\vartheta=1$

$*10^{-5}$

Upper CLEO bound

SM

$\tan\beta=20$

Lower CLEO bound

BR(b→sγ)

M_g(GeV)

The BR vs gluino mass for $\mu<0$

$M_{tpole}=178GeV$ $\sin\vartheta=1$

$\tan\beta=23$

$*10^{-5}$

Upper CLEO bound

SM

Lower CLEO bound

BR(b→sγ)

M_g(GeV)

The BR vs gluino mass for $\mu > 0$

$M_{tpole} = 178 GeV$

$*10^{-4}$

BR($b \rightarrow s\gamma$)

20

15

10

2

Current CLEO upper bound

M_g(GeV)

The BR vs gluino mass for $\mu<0$

$\vartheta=(2\pi)/3$ $M_{tpole}=178GeV$

$*10^{-5}$

Upper CLEO bound

SM

6

10

15

Lower CLEO bound

BR($b\to s\gamma$)

M_g(GeV)

The BR vs gluino mass for $\mu>0$

The BR vs gluino mass for $\mu > 0$